

# Neutrino-photon reactions for energies above $m_e$

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## Abstract

We show that neutrino-photon reactions above  $m_e$  are dominated by the reaction  $\nu\gamma \rightarrow \nu e^+ e^-$ . We calculate its cross-section and see that it is larger by several orders of magnitude than the cross-sections of other neutrino-photon processes, for energies above  $m_e$ . We also discuss potential astrophysical and cosmological consequences.

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Neutrino-photon reactions could play a role in some astrophysical or cosmological environments. Recently, there has been some interest in these type of interactions. (See Refs. [1]-[9].) Let us consider

$$\nu \gamma \rightarrow \nu X \quad (1)$$

This reaction has to involve weak interactions since the neutrino has no electric charge and the expected magnetic moment is small. The elastic channel

$$\nu \gamma \rightarrow \nu \gamma \quad (2)$$

is further suppressed due to the prohibition of a two-photon coupling to a  $J = 1$  state (Yang theorem). The cross-section is on the order of  $\sigma \sim G_F^2 \alpha^2 \omega^6 / M_W^4$  [1, 2] where  $\omega$  is the energy of the photon in the center-of-mass system ( $\sqrt{s} = 2\omega$ ) and where we assume we are in the limit  $\omega \ll M_W$  and also  $m_\nu \ll \omega$ . The Yang suppression amounts to the small factor  $(\omega/M_W)^4$  in  $\sigma$ . We refer the reader to [2] where the cross-section of the process (2) in the standard model has been calculated numerically.

Recently, in a series of papers [3]-[6] it has been realized that the inelastic process

$$\nu \gamma \rightarrow \nu \gamma \gamma \quad (3)$$

largely dominates over (2). The technicalities for the calculation of the cross-section are different depending on whether we have  $\omega \ll m_e$  or not.

For  $\omega \ll m_e$ , one has a cross-section of magnitude  $\sigma \sim G_F^2 \alpha^3 \omega^{10} / m_e^8$  [3, 5]. Now there is no Yang suppression, as can be readily seen in the expression for the cross-section, and instead we have an inverse power of the electron mass. Indeed, (3) is related to  $\gamma \gamma \rightarrow \gamma \gamma$  scattering when substituting a photon by the neutrino current, and for  $\omega \ll m_e$  one can make use of the effective Euler-Heisenberg lagrangian that has the inverse power  $m_e^{-4}$  of the scale  $m_e$ . This simplifies the calculation and explains the appearance of the electron mass in the cross-section. The change of the  $M_W$  scale by the  $m_e$  scale explains the enhancement of the inelastic channel (3).

When  $\omega > m_e$ , the cross-section has to be calculated directly and there is no simplification of the type we have mentioned. The calculation (for  $\omega \ll M_W$ ) has been carried in [4, 6]. From a theoretical point of view, this calculation is interesting since it is a nice example where one can compare the results using an effective lagrangian and the results of a full calculation that can be used at high energies. The authors argue that the region  $\omega > m_e$

may have some phenomenological interest. For example, if one wants to check whether the reaction (3) contributes to the neutrino opacities in the first stages of a supernova one needs clearly the cross-section when  $\omega > m_e$ . Also, the authors of [4] as well as the authors of [6] consider some possible cosmological applications. They notice that, if there is any of such applications, it would occur in the regime  $\omega > m_e$ .

We would like to point out that when one crosses the  $m_e$  threshold the reaction

$$\nu \gamma \rightarrow \nu e^+ e^- \quad (4)$$

dominates over the reaction (3). In fact, the cross-sections are larger by several orders of magnitude than the cross-sections of (3). The reason is simple to understand if one considers the limit  $M_W \gg \omega \gg m_e$ . In this limit, while (3) is on the order of  $G_F^2 \alpha^3 \omega^2$ , (4) is of order  $G_F^2 \alpha \omega^2$ . Due to the potential interest of the process, we would like to present its calculation for all energies  $\omega$  (below  $M_W$ ).

First of all, we show the associated Feynman diagrams for the process in figure 1. We assume the limit  $\omega \ll M_W$ , so that we can use the Fermi effective coupling,

$$\mathcal{L}_F = -\frac{4G_F}{\sqrt{2}} \bar{\Psi}_\nu \gamma_\mu P_L \Psi_\nu \bar{\Psi}_e \gamma^\mu (c_L P_L + c_R P_R) \Psi_e \quad (5)$$

with  $P_L = (1 - \gamma_5)/2$  and  $P_R = (1 + \gamma_5)/2$ . When  $\nu = \nu_\mu$  or  $\nu = \nu_\tau$  there is only the neutral weak current contributing to (5), and we have

$$\begin{aligned} c_L &= -1/2 + \sin^2 \theta_W \\ c_R &= \sin^2 \theta_W \end{aligned} \quad (6)$$

( $\theta_W$  is the weak mixing angle). When  $\nu = \nu_e$ , apart from the neutral current there is also the charged current contribution. It can be Fierz rearranged to have the form of (5), and then

$$\begin{aligned} c_L &= 1/2 + \sin^2 \theta_W \\ c_R &= \sin^2 \theta_W \end{aligned} \quad (7)$$

Next, we write the amplitude for  $\nu(p) + \gamma(k) \rightarrow \nu(q) + e^+(q_1) + e^-(q_2)$

$$\begin{aligned} M &= -\frac{4ieG_F}{\sqrt{2}} \epsilon_\mu(k) \bar{u}(q) \gamma_\nu P_L u(p) \times \\ &\bar{u}(q_1) \left[ \gamma^\mu \frac{1}{\not{q}_1 - \not{k} - m} \gamma^\nu (c_L P_L + c_R P_R) + \gamma^\nu (c_L P_L + c_R P_R) \frac{1}{-\not{q}_2 + \not{k} - m} \gamma^\mu \right] v(q_2) \end{aligned} \quad (8)$$

Finally, we evaluate numerically the cross-section. We work, as we said, in the limit  $\omega \ll M_W$ , as well as  $m_\nu \ll \omega$ . The results, for the case  $\nu = \nu_e$  are presented in figure 2, in the range  $m_e < \omega < 10^2 m_e$ . For higher energies, the cross-section scales approximately as  $\omega^2$ . For instance, for  $\omega = 100$  MeV, we have  $\sigma \simeq 2 \times 10^{-3}$  fb, and for  $\omega = 1$  GeV, it is already  $\sigma \simeq 0.4$  fb. The cross-section for our process (4) when  $\nu$  is either  $\nu_\mu$  or  $\nu_\tau$  is obtained when using (6) instead of (7). As expected, one gets numerically a similar result.

In figure 2 we also show the cross-section for the reaction (3), and as expected we see that it is smaller by several orders of magnitude. Of course, our reaction has a threshold at  $\omega = m_e$ . For energies below  $m_e$ , the dominant process is (3).

Let us now examine potential consequences in the supernova dynamics. Any reaction of the type (1) contributes to the neutrino opacity in a supernova collapse. In the conservative range of temperatures  $T = 10 - 100$  MeV, both photons and neutrinos have energies exceeding  $m_e$  and thus it is our reaction (4) that will be most important. To evaluate its possible role, we need the neutrino mean free path,

$$\lambda = \frac{1}{\sigma n_\gamma} \quad (9)$$

due to (4). We estimate  $\sigma$  by evaluating it at the average energy  $\omega \sim 3T$ . In the temperature range we have indicated, we get

$$\lambda \sim 3 \times 10^8 - 2 \times 10^3 \text{ cm} \quad (10)$$

(the shorter  $\lambda$  corresponding to the higher temperature, of course).

Neutrino scattering with non-relativistic nucleons in a supernova has a cross-section  $\sigma \sim 2 \times 10^{-5} (E/\text{MeV})^2$  fb, where  $E$  is the neutrino energy. The large nucleon density,  $n_N \sim 2 \times 10^{38} \text{ cm}^{-3}$ , leads to  $\lambda = 1/\sigma n_N \sim 3 - 300$  cm. At the view of the figures in (10) we tentatively conclude that the role of neutrino-photon reactions in the neutrino opacity in the supernova is small.

Regarding cosmological applications, the interest of reactions (1) in the early universe has been discussed in Refs. [7] and [6] (See also Ref.[8].) In the standard scenario, we have that for temperatures  $T > 1$  MeV, weak interactions have not yet decoupled and neutrinos interact with matter, in particular with electrons. Electrons in turn interact with photons. Thus, neutrinos, photons and electrons are in equilibrium. At  $T \sim 1$  MeV, neutrinos decouple.

Once we have direct neutrino-photon interactions, it may be interesting to investigate at which temperature  $T_D$  these direct interactions decouple. In [7] it is shown that, due to the reaction (2), neutrinos decouple from photons at  $T_D \sim 1$  GeV. In [6], the authors consider the inelastic process (3) and demonstrate that  $T_D \sim 1$  GeV. We expect a lower decoupling temperature due to our reaction (4). Let us show that this is the case.

We have, on the one hand that the expansion rate of the Universe is given by

$$H = 1.66 \, g^{1/2} \frac{T^2}{M_{Pl}} \quad (11)$$

where  $M_{Pl}$  is the Plank energy. Anticipating that  $1 \text{ MeV} < T_D < 100 \text{ MeV}$ , we set the total degrees of freedom  $g = 10.75$ . On the other hand, the interaction rate is given by

$$\Gamma = \sigma n_\gamma \quad (12)$$

Again, to estimate  $\sigma$ , we use the average energy  $\omega \sim 3T$ . At high energies, one has that the ratio  $\Gamma/H \gg 1$ , as expected. This ratio decreases with decreasing temperature and becomes of order unity for  $T = T_D \sim 10 \text{ MeV}$ . This is the decoupling temperature of neutrinos and photons due to the process (4), and we see that is below the one found using other neutrino-photon reactions. In any case, neutrinos are not decoupled due to weak interactions until  $T_D \sim 1 \text{ MeV}$ , and we are not able to see any interesting application of our result.

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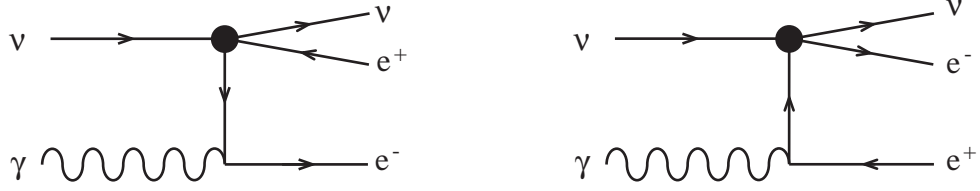


FIG. 1: Feynman diagrams for the process  $\nu\gamma \rightarrow \nu e^+ e^-$ .

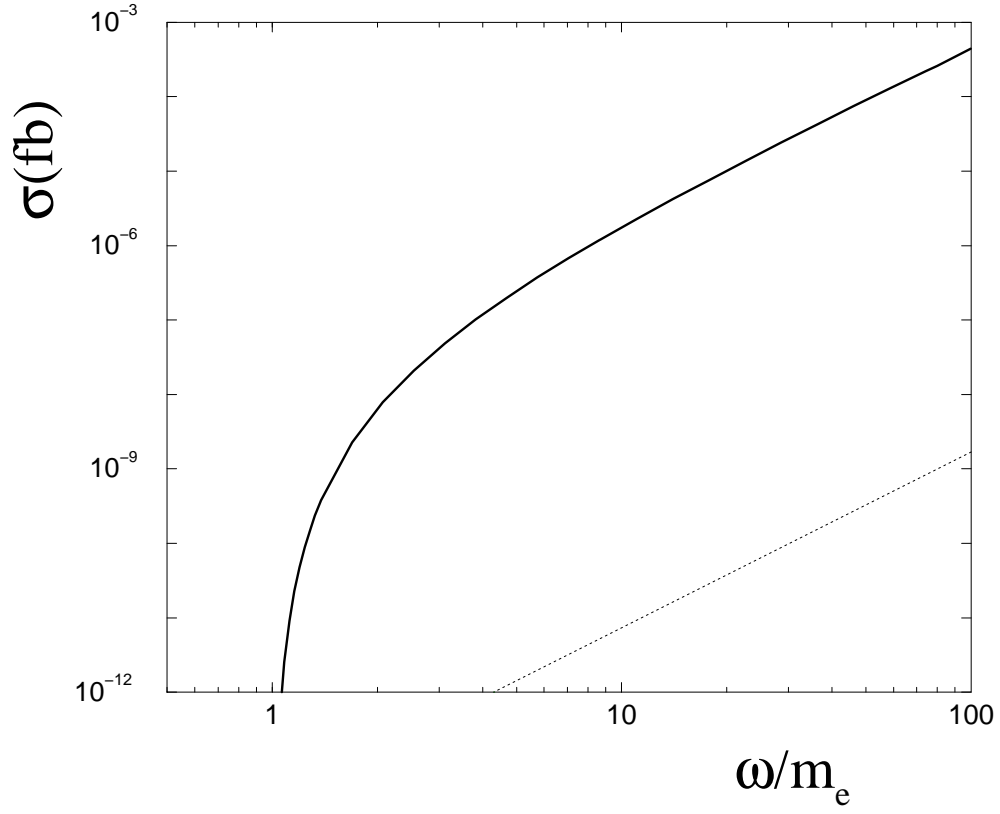


FIG. 2: The cross-section for the process  $\nu\gamma \rightarrow \nu e^+ e^-$  in fb as a function of  $\omega/m_e$  with  $\omega$  the center-of-mass energy of the photon. We display the cross-section for the process  $\nu\gamma \rightarrow \nu\gamma\gamma$  (dotted line).